

A Singlet-Extension of the MSSM for 125 GeV Higgs with the Least Tuning

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Abstract

In order to raise the Higgs mass up to 125 GeV and relieve the fine-tuning associated with the heavy s-top mass in the MSSM, we propose a new singlet extension of the MSSM. In this scenario, the Higgs mass is radiatively generated also in a hidden sector, and the effect is transmitted to the Higgs through a messenger field. The Higgs mass can be efficiently raised by the parameters of the superpotential as in the extra matter scenario, but free from the constraints on extra colored matter fields by the LHC experiments. The tuning problem can be remarkably mitigated by taking low enough messenger (~ 300 GeV) and mass parameter scales (~ 500 GeV).

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I. INTRODUCTION

Recently, CMS and ATLAS reported the observations of the signals, which can be interpreted as the presence of a standard model(-like) Higgs with the mass of 125 GeV around the five sigma confidence level [1, 2]. The news seems to be accepted as the discovery of the long-awaited Higgs particle, which is very essential in mass generations for the standard model (SM) particles. However, the theoretical issues associated with the Higgs boson, e.g. how the Higgs can naturally exist at low energies, still remain unsolved. Actually, they have played the role of strong motivations to study various new physics beyond the SM.

For the last three decades, the minimal supersymmetric standard model (MSSM) has maintained the status of the leading candidate beyond the SM [3]. It provides a beautiful solution to the large hierarchy problem between the electroweak (EW) energy scale and the grand unification (GUT) or Planck scale with the minimal extension of the SM in the supersymmetric (SUSY) manner. It makes it possible to embed the SM in a fundamental theory like the string theory [4]. The gauge coupling unification is another great advantage of the MSSM.

MSSM: In the MSSM, a relatively lower value of the Higgs mass is preferred. It is basically because the tree-level quartic coupling of the Higgs potential is given by the small gauge coupling unlike the SM. As a result, the Higgs mass cannot be larger than even the Z boson mass (M_Z), were it not for the large radiative correction: by including the radiative correction by the large top quark Yukawa coupling, the Higgs mass can be lifted above 100 GeV. Actually, the lightest Higgs mass in the MSSM is given by

$$m_h^2 \approx M_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} (y_t m_t)^2 \sin^2 \beta \log \left(\frac{m_t^2 + \tilde{m}_t^2}{m_t^2} \right), \quad (1)$$

where y_t is the top quark Yukawa coupling, and m_t^2 and \tilde{m}_t^2 denote the mass squared of the top quark and the soft mass squared of its superpartner “s-top,” respectively. The first term in the right hand side comes from the tree level contribution, and the second term from the radiative correction ($\equiv \Delta m_h^2|_{\text{MSSM}}$) by the top quark and the s-top. Here we neglect the “A-term” contribution. In Eq. (1), however, the values of the top quark mass m_t and also the top quark Yukawa coupling y_t have been already precisely measured. Thus, only a useful parameter for raising the Higgs mass is the soft mass squared of the s-top. Note that the radiative correction *logarithmically* depends on the s-quark mass squared. Thus, raising the Higgs mass with the soft mass squared of the s-top is not a quite efficient way. Indeed, a s-top mass heavier than a few TeV is needed to achieve 125 GeV Higgs mass at two-loop level, unless the large mixing effect between the left and right s-tops through the A-term contribution assumed [3, 5]. However, the s-top mass cannot be arbitrarily large.

The radiative correction by the top and s-top contributes also to the renormalization of

the soft parameter m_2^2 , which is the soft mass squared of the MSSM Higgs “ H_u ,”

$$m_2^2(M_Z) \approx m_2^2 - \frac{3y_t^2}{8\pi^2} \tilde{m}_t^2 \log \left(\frac{M_G^2}{m_t^2 + \tilde{m}_t^2} \right). \quad (2)$$

Here we keep only the radiative correction coming from the top quark Yukawa coupling, which is the largest correction to m_2^2 . Here M_G indicates the GUT scale ($\approx 2 \times 10^{16}$ GeV), at which the soft parameters are assumed to be generated in the minimal supergravity model. The negative contribution of it causes flipping the sign of m_2^2 at the EW energy scale, which triggers the EW symmetry breaking. Thus, one of the extremum conditions for the MSSM Higgs fields is modified as

$$m_2^2 + |\mu|^2 \approx m_3^2 \cot \beta + \frac{M_Z^2}{2} \cos 2\beta - \frac{3y_t^2}{8\pi^2} \tilde{m}_t^2 \log \left(\frac{m_t^2 + \tilde{m}_t^2}{M_G^2} \right). \quad (3)$$

The radiative corrections add the last term ($\equiv \Delta m_2^2$) in Eq. (3). If a too heavy s-top mass is taken to raise the Higgs mass from Eq. (1), Δm_2^2 and other parameters should be properly tuned to give M_Z^2 , which implies that the EW symmetry breaking becomes unnatural. Actually, Eq. (3) is not directly related to the observed value of the Higgs mass, but is closely associated with the naturalness of the EW symmetry breaking. It is known as the name of the “little hierarchy problem” in the MSSM. Thus, e.g. for \tilde{m}_t of 2 TeV, the size of the tuning is roughly estimated with the hierarchy in the relation of Eq. (3):

$$\frac{(M_Z^2/2) \cos 2\beta}{|\Delta m_2^2|} < \left| \left(\frac{\tilde{m}_t^2}{M_Z^2} \right) \frac{3y_t^2}{4\pi^2} \log \left(\frac{\tilde{m}_t^2}{M_G^2} \right) \right|^{-1} \lesssim 4.7 \times 10^{-4}. \quad (4)$$

In order to reduce the tuning in Eq. (3), thus,

- smaller mass parameters need to be taken, but yielding $m_h = 125$ GeV, and
- a low energy soft term generation scenario is needed for a smaller log piece in Eq. (3).

In this paper, we will introduce a phenomenologically attractive scenario, addressing the above two requirements.

Maximal Mixing: In fact, 125 GeV Higgs mass could be achieved even with relatively lighter s-tops by considering also the “A-term” contribution to the radiative correction, which was dropped in Eq. (1). A large mixing between the s-tops of the $SU(2)_L$ doublet and singlet, $(\tilde{t}_L, \tilde{t}_R)$, via the SUSY breaking “A-term” is very helpful for raising the Higgs mass. Particularly, the “maximal mixing”

$$X_t \equiv (A_t - \mu \cot \beta) = \sqrt{6} m_{\tilde{t}}, \quad (5)$$

where $m_{\tilde{t}} \equiv \sqrt{m_t^2 + \tilde{m}_t^2}$, can lift the Higgs mass up to 135 GeV without any other helps in the decoupling limit of the CP odd Higgs [3]. However, as the mixing deviated from the maximal mixing, the enhancement effect drops rapidly. Employing a large mixing of \tilde{t}_L - \tilde{t}_R ,

hence, would be a kind of fine-tuning in this sense. Throughout this paper, we will not consider such a mixing effect.

Extra Matter: In order to efficiently enhance the radiative correction, one might introduce the fourth family of the chiral matter or extra vector-like matter [6, 7]. In the case of the fourth family of the chiral matter, the top quark Yukawa coupling and also the top quark mass in Eq. (1) are replaced by the unknown parameters, which can be utilized to raise the Higgs mass. Since such SUSY parameters appears outside the logarithmic function, they can efficiently increase the Higgs mass unlike the s-top mass squared in the MSSM. Unfortunately, however, the presence of extra colored particles coupled with the Higgs with order one Yukawa couplings would exceedingly affect the production rate and also decay rate of the Higgs at the large hadron collider (LHC) i.e. $gg \rightarrow h$ and $h \rightarrow \gamma\gamma$: they result in immoderate deviation from the LHC data. According to Ref. [8], indeed, the existence of such a extra family of the chiral matter is excluded at the 99.9% confidence level for the Higgs mass of 125 GeV.

In the case of extra vector-like matter, in which a Yukawa coupling of order unity with the Higgs is still necessary for lifting the Higgs mass, the LHC bound could be avoided by employing heavy enough mass terms for vector-like fields. However, the tuning problem associated with the naturalness of the Higgs mass becomes serious with the high scale mass parameters.¹ Moreover, the extra vector-like matter should compose the SU(5) or SO(10) multi-plets to protect the gauge coupling unification. If the low energy effective theory is not embedded in four dimensional SU(5) or SO(10) GUTs but in other unified theory defined in higher dimensional spacetime like string theory [4], we need to explore other possibilities for explaining the 125 GeV Higgs mass.

NMSSM: In the next-to-minimal supersymmetric standard model (NMSSM), the Higgs mass can be raised using the tree level correction of the Higgs potential [9–11]. In the NMSSM, the MSSM μ term is promoted to a renormalizable trilinear term SH_uH_d in the superpotential, introducing an extra singlet superfield S together with a dimensionless coupling λ . The presence of such a trilinear term in the superpotential provides the quartic coupling in the Higgs potential as well as a μ problem solution through the gravity mediated SUSY breaking scenario. By the quartic Higgs potential coming from λSH_uH_d in the

¹ For instance, if only an extra vector-like pair of quark doublets $\{Q, Q^c\}$ is introduced and the superpotential $W = M_Q QQ^c + yQH_u u^c$, where H_u and u^c are the Higgs and a quark singlet in the MSSM, is considered, using the formula in [7] one can show that the radiative correction to the Higgs potential is

$$\Delta V = \frac{3}{16\pi^2} \left[(M^2 + \tilde{m}^2)^2 \left\{ \log \left(\frac{M^2 + \tilde{m}^2}{\Lambda^2} \right) - \frac{3}{2} \right\} - M^4 \left\{ \log \left(\frac{M^2}{\Lambda^2} \right) - \frac{3}{2} \right\} \right] + \text{constant}, \quad (6)$$

where $M^2 \equiv M_Q^2 + y^2 |H_u|^2$ and Λ indicates a renormalization scale. Here all the soft mass squareds are set to be \tilde{m}^2 , and the “A-term” effect is ignored for simplicity. This expression is quite similar to that in the case of Ref. [12]. However, the fields circulating on the loops in Ref. [12] are MSSM singlets.

superpotential, the mass of the lighter CP even Higgs in the NMSSM is modified at tree level as

$$m_h^2 \approx M_Z^2 \cos^2 2\beta + \lambda^2 v_H^2 \sin^2 2\beta + \Delta m_h^2|_{\text{MSSM}} , \quad (7)$$

where $v_H^2 \equiv v_u^2 + v_d^2 = (174 \text{ GeV})^2$, and $\Delta m_h^2|_{\text{MSSM}}$ denotes the radiative correction by the (s-) top. The tree level correction “ $\lambda^2 v_H^2 \sin^2 2\beta$ ” in Eq. (7) can remarkably raise the Higgs mass, if the dimensionless Yukawa coupling λ was sizable. In order to maintain the perturbativity of the model up to the GUT scale, however, λ is known to be smaller than 0.7 at the EW scale (“Landau pole constraint”) [9]. Moreover, to achieve the Higgs mass of 125 GeV with the s-top mass much lighter than 1 TeV, which is necessary for the naturalness of the Higgs, λ needs to be larger than 0.5. Requiring both the perturbativity and the naturalness, thus, the range of λ should be quite limited:

$$0.5 \lesssim \lambda \lesssim 0.7. \quad (8)$$

The relatively small λ pushes $\tan\beta$ to the smaller values for obtaining 125 GeV Higgs mass,

$$1 \lesssim \tan\beta \lesssim 3, \quad (9)$$

which gives almost the maximal values to $\sin^2 2\beta$ in Eq. (7).

Radiative Correction by MSSM Singlets: Recently, the authors of Ref. [12] proposed a scenario in which the Higgs mass is raised through *radiative corrections by some MSSM singlet fields*. In this case, the Higgs mass can be efficiently lifted by using the parameters of the superpotential just like the extra matter case, but the LHC constraint can be avoided because only MSSM singlets are employed. In Ref. [12], it was shown that the parameter space of $\tan\beta$ and the trilinear coupling of “ $SH_u H_d$ ” ($\equiv y_H$) in the superpotential for explaining the 125 GeV Higgs mass can be remarkably enlarged by extending the NMSSM with some other MSSM singlets, compared to the original form of the NMSSM: $0.2 \lesssim y_H \lesssim 0.5$ and $3 \lesssim \tan\beta \lesssim 10$ can be also consistent with the Higgs mass of 125 GeV even without considering the mixing effect.

Since the Higgs mass is radiatively generated from a hidden sector and then it is transmitted to the Higgs sector through a mediation by a messenger in this scenario, the fine-tuning problem can be quite alleviated by taking low scale messenger and mass parameters. In this paper, we will particularly discuss how much the fine-tuning in the Higgs sector can be relieved in this setup.

This paper is organized as follows. In section II, our basic setup will be introduced. In section III, the effective Higgs potential will be calculated in our setup. In section IV, we will discuss how to achieve the 125 GeV Higgs mass and minimize the tuning. In section V, we will propose a UV model. Section VI will be devoted to conclusion.

II. A SINGLET-EXTENSION OF THE MSSM

In this paper, we will pursue the naturalness of the model rather than its minimality. Introducing the *MSSM singlet* superfields $\{S, S^c\}$ and $\{N, N^c\}$, we extend the MSSM Higgs sector in the superpotential as follows:

$$W = (\mu + y_H S) H_u H_d + \mu_S S S^c + (\mu_N + y_N S^c) N N^c \quad (10)$$

where $\{H_u, H_d\}$ denote the two MSSM Higgs doublets. For simplicity we assume the parameters in Eq. (10) are all real. Since the μ and $\mu_{S,N}$ terms are explicitly present, there does not remain a Pecci-Quinn (PQ) symmetry at the EW scale. Apart from the MSSM μ term, the trilinear term $SH_u H_d$ *a la* the NMSSM is introduced in Eq. (10) [13]. Actually, Eq. (10) should be regarded as a low energy effective superpotential. How such a superpotential can be generated will be discussed later.

$\{S, S^c\}$ are the messenger fields, which connect the Higgs $\{H_u, H_d\}$ and the hidden sector fields $\{N, N^c\}$. Note that the “messenger” and “hidden sector” here do not necessarily mean the conventional ones appearing in various SUSY breaking scenarios. We only require the mass splitting in the hidden sector superfields $\{N, N^c\}$ such that they eventually generate the radiative correction of the Higgs mass. Such an effect can be transmitted to the Higgs via the messengers $\{\tilde{S}, \tilde{S}^c\}$ as will be seen later. $\{N, N^c\}$ are a vector-like n -dimensional representation of a certain hidden gauge group. They could remain light down to low energies by the global symmetries.

$\mu_{S,N}$ terms are the Dirac type bare mass terms of the messengers and hidden sector fields. $\mu_{S,N}$ are assumed to be larger than 300 GeV. Thus, the squared masses of $\{\tilde{S}, \tilde{S}^c\}$ and $\{\tilde{N}, \tilde{N}^c\}$ are quite heavier than that of the lightest Higgs. Since μ_S and μ_N both are much heavier than the Higgs mass, there is no “singlet-ino” (the fermionic components of singlet superfields) lighter than the Higgs. Hence, there is no invisible decay channel of the Higgs in this model. However, we restrict $\mu_{S,N}$ to be smaller than 1 TeV. It is because the fine-tuning in the Higgs sector would become serious, if they are heavier than 1 TeV. Their smallness compared to the fundamental scale will be explained in section V.

In fact, the superpotential Eq. (10) can provide a quartic Higgs potential at tree level via $|\partial W / \partial S|^2$ as in the NMSSM, which is quite helpful for lifting the Higgs mass if y_H was sizable. However, the Landau pole constraint to avoid blowing-up below the GUT scale is known to restrict the size of y_H to be smaller than 0.7 [9]. While y_H should be smaller than the unity, y_N *can be still of order unity* at the EW scale. Nonetheless, the hidden gauge interaction of $\{N, N^c\}$ can prevent y_N from blowing-up at higher energy scales, because $\{N, N^c\}$ carry a non-Abelian gauge charge of a relatively large hidden gauge group.

Since y_H is relatively small and μ_S is quite heavier than the Higgs mass, the tree level correction by $\{S, S^c\}$ to the Higgs potential is expected to be suppressed. Moreover, the mixing angles between the Higgs and the singlet sectors would be negligible. In Ref. [12], however, it was shown that even with relatively small y_H (0.2-0.5), the Higgs mass of 125 GeV

can be achieved through the large radiative correction, if a relatively larger y_N compensates the smallness of y_H .

With small enough y_H the soft mass squared of S , \tilde{m}_S^2 does not run much with energy at one-loop level. On the other hand, y_N is of order unity, and so $\tilde{m}_{S^c}^2$ can be suppressed at low energies compared to \tilde{m}_S^2 by the renormalization group (RG) effect. Due to the gauge interaction in the hidden sector, the soft masses of N and N^c , \tilde{m}_N and \tilde{m}_{N^c} can be quite heavier than other soft masses at low energies. For simplicity of the future calculation, but considering the RG behaviors, we assume a hierarchy among the mass parameters:

$$\tilde{m}_{S^c} \lesssim \mu \lesssim m_{3/2}, \mu_S \lesssim \tilde{m}_S \lesssim \mu_N, \tilde{m}_N (= \tilde{m}_{N^c}), \quad (11)$$

where $m_{3/2}$ collectively denotes typical soft parameters except \tilde{m}_S and \tilde{m}_{S^c} . Although \tilde{m}_{S^c} is the smallest, the scalar component of S^c is still much heavier than the Higgs because its physical mass squared is given by $\mu_S^2 + \tilde{m}_{S^c}^2$.

III. THE EFFECTIVE HIGGS POTENTIAL

Let us first integrate out the quantum fluctuations of $\{N, N^c\}$. Due to the mass difference between the bosonic and fermionic components in $\{N, N^c\}$, the one-loop effective potential of \tilde{S}^c is generated [14]:

$$\Delta V = \frac{n}{16\pi^2} \left[(M_N^2 + \tilde{m}_N^2)^2 \left\{ \log \left(\frac{M_N^2 + \tilde{m}_N^2}{\Lambda^2} \right) - \frac{3}{2} \right\} - M_N^4 \left\{ \log \left(\frac{M_N^2}{\Lambda^2} \right) - \frac{3}{2} \right\} \right], \quad (12)$$

where Λ denotes a renormalization mass scale. The SUSY mass of $\{N, N^c\}$ ($\equiv M_N$) is given by the summation of μ_N and the classical value of \tilde{S}^c as explicitly seen in the superpotential Eq. (10), and so

$$M_N^2 = \left| \mu_N + y_N \tilde{S}^c \right|^2. \quad (13)$$

Thus, ΔV in Eq. (12) depends only on \tilde{S}^c .

Including the soft terms and the one-loop effective potential obtained after integrating out $\{N, N^c\}$, $\Delta V(\tilde{S}^c)$, the scalar potential associated with the superpotential Eq. (10) is derived as follows:

$$\begin{aligned} V_{\text{HS}} = & \left(m_2^2 + |\mu + y_H \tilde{S}|^2 \right) |H_u|^2 + \left(m_1^2 + |\mu + y_H \tilde{S}|^2 \right) |H_d|^2 \\ & + \left(\tilde{m}_{S^c}^2 + \mu_S^2 \right) |\tilde{S}^c|^2 + \left(\tilde{m}_S^2 + \mu_S^2 \right) |\tilde{S}|^2 + y_H^2 |H_u H_d|^2 \\ & + \left[\left(y_H \mu_S \tilde{S}^{c*} + B_\mu \mu + y_H A_S \tilde{S} \right) H_u H_d + B_S \mu_S \tilde{S} \tilde{S}^c + \text{h.c.} \right] \\ & + \frac{1}{8} (g^2 + g'^2) (|H_u|^2 - |H_d|^2)^2 + \frac{1}{2} g^2 |H_u^\dagger H_d|^2 + \Delta V(\tilde{S}^c), \end{aligned} \quad (14)$$

where we set $\tilde{N} = \tilde{N}^c = 0$, since the vanishing field value conditions for such heavy scalars fulfill all the extremum conditions of the scalar potential.

Now let us integrate out $\{\tilde{S}, \tilde{S}^c\}$, which are heavier than $\{H_u, H_d\}$. The equations of motion in the static limit for $\{\tilde{S}, \tilde{S}^c\}$ are

$$\begin{aligned}\frac{\partial V_{\text{HS}}}{\partial \tilde{S}^c} &= (\tilde{m}_{S^c}^2 + \mu_S^2) \tilde{S}^{c*} + B_S \mu_S \tilde{S} + y_H \mu_S H_u^* H_d^* + \partial_{\tilde{S}^c} \Delta V = 0, \\ \frac{\partial V_{\text{HS}}}{\partial \tilde{S}} &= (\tilde{m}_S^2 + \mu_S^2) \tilde{S}^* + B_S \mu_S \tilde{S}^c + y_H A_S H_u H_d + (y_H \mu + y_H^2 \tilde{S}^*) (|H_u|^2 + |H_d|^2) = 0.\end{aligned}\quad (15)$$

Considering the hierarchy suggested in Eq. (11), the approximate solutions to Eq. (15) are given by

$$\begin{aligned}\tilde{S}^c &\approx \frac{-1}{\mu_S^2} \left[y_H \mu_S H_u H_d (1 + \epsilon_1 - \epsilon_2^*) + \partial_{\tilde{S}^{c*}} \Delta V^* (1 + \epsilon_1) \right], \\ \tilde{S} &\approx \frac{-1}{\tilde{m}_S^2 + \mu_S^2} \left[y_H (A_S^* - B_S^*) H_u^* H_d^* - \frac{B_S^*}{\mu_S} \partial_{\tilde{S}^c} \Delta V \right] \ll \tilde{S}^c,\end{aligned}\quad (16)$$

where the terms proportional to \tilde{m}_{S^c} and μ are ignored due to their relative smallness in Eq. (11), and ϵ_1 and ϵ_2 are defined as

$$\epsilon_1 \equiv \frac{|B_S|^2}{\tilde{m}_S^2 + \mu_S^2} \quad \text{and} \quad \epsilon_2 \equiv \frac{A_S^* B_S}{\tilde{m}_S^2 + \mu_S^2}, \quad (17)$$

respectively. Inserting the expressions of the heavy fields in Eq. (16) into the scalar potential V_{HS} of Eq. (14) yields the low energy effective Higgs potential:

$$\begin{aligned}V_H &\approx (m_2^2 + \mu^2) |H_u|^2 + (m_1^2 + \mu^2) |H_d|^2 + (B_\mu \mu H_u H_d + \text{h.c.}) \\ &\quad + \frac{1}{8} (g^2 + g'^2) (|H_u|^2 - |H_d|^2)^2 + \frac{1}{2} g^2 |H_u^\dagger H_d|^2 \\ &\quad - \frac{|A_S - B_S|^2}{\tilde{m}_S^2 + \mu_S^2} y_H^2 |H_u H_d|^2 + \Delta V(H).\end{aligned}\quad (18)$$

Here we dropped the (reducible) 2-loop effects coming from $|\partial_{\tilde{S}^c} \Delta V|^2$. Note that the first two lines in Eq. (18) are nothing but the MSSM Higgs potential, while the two terms in the third line correspond to the tree-level and one-loop corrections induced by the heavy fields $\{\tilde{S}, \tilde{S}^c\}$ and $\{N, N^c\}$. The quartic term “ $y_H^2 |H_u H_d|^2$ ” in Eq. (14) is cancelled out, and so as seen from Eq. (18), the tree-level correction remains quite suppressed with heavy mass parameter \tilde{m}_S^2 and μ_S^2 . As will be seen later, however, the one-loop correction $\Delta V(H)$ can be relatively large, since it is originated from other sector rather than the MSSM.

The one-loop correction $\Delta V(H)$ in Eq. (18) is just given by Eq. (12), but the M_N in its expression should be replaced by

$$M_N^2 \approx \mu_N^2 - \left(y_H y_N \frac{\mu_N}{\mu_S} H_u H_d + \text{h.c.} \right) \equiv \mu_N^2 - \left(y_H y_N \frac{\mu_N}{\mu_S} \right) h_u h_d, \quad (19)$$

using Eq. (16). Here $h_{u,d}$ is the real component of $H_{u,d}$, $\text{Re} H_{u,d} \equiv \frac{1}{\sqrt{2}} h_{u,d}$. Thus, the expression of $\Delta V(H)$ here is exactly the same as that of Ref. [12]. In Ref. [12], $\{N, N^c\}$

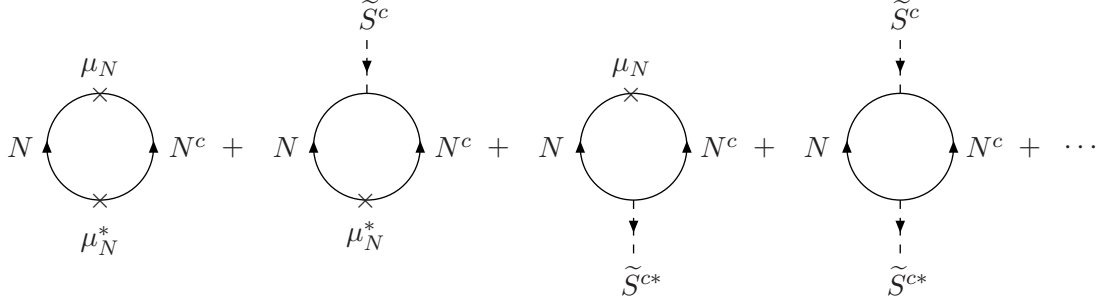


FIG. 1: Some contributions to the one-loop effective potential of \tilde{S}^c . Here we present only the diagrams of the fermionic loops. By infinite summation of the amplitudes for all the relevant one-loop diagrams, the Coleman-Weinberg's effective potential for \tilde{S}^c can be derived. Below the mass scale of \tilde{S}^c , the low energy effective Higgs potential can be obtained by integrating out \tilde{S}^c , in which “ $H_u H_d$ ” is attached to \tilde{S}^c in this setup.

are integrated out after $\{S, S^c\}$. As pointed out in Ref. [12], however, the result should be insensitive to the sequence of the decouplings, since the mass scales of $\{N, N^c\}$ and $\{S, S^c\}$ are not much hierarchical.

We note that the similarity between the one-loop effective potential of Eq.(12) with Eq. (19) and that of the footnote 1 in Introduction, which is the radiative Higgs potential for a simple case of extra vector-like matter. Accordingly, one can expect that the Higgs mass is raised in our case in a similar way to the case of extra vector-like matter. The most important difference between these two scenarios is that the fields circulating on the loops are MSSM singlets in our case, while they are charged fields under the SM in the extra vector-like matter case. In our case, lower scale mass parameters can be taken for e.g. alleviating the tuning problem, but the LHC constraint on the extra colored particles can be avoided unlike the extra vector-like matter case.

In fact, the Coleman-Weinberg's one-loop effective potential, $\Delta V(\tilde{S}^c)$ of Eq. (12) with Eq. (13), can be obtained by taking infinite summation of all possible one-loop diagrams, in which arbitrary numbers of \tilde{S}^c are attached on the loop as the external legs [14]. See the diagrams of FIG. 1, in which only the diagrams of the fermionic loops are presented. In the effective operators valid below the mass scale of \tilde{S}^c , however, \tilde{S}^c should appear as internal legs. As seen from Eq. (14), \tilde{S}^{c*} interacts only with $H_u H_d$ at tree level, the external legs of \tilde{S}^c in FIG.1 can be hidden only with $H_u H_d$ at one-loop level. See FIG. 2-(b). In fact, \tilde{S}^{c*} and \tilde{S} are mixed and \tilde{S} is also coupled to $H_u H_d$ via the B_S and A_s terms. Thus, \tilde{S}^c can couple to $H_u H_d$ through \tilde{S} . However, this possibility is more suppressed due to the hierarchical mass relation in Eq. (11).

In this scenario, a non-zero radiative correction to the Higgs mass squared is generated by the mass splitting of $\{N, N^c\}$ in the hidden sector. The hidden sector in this model,

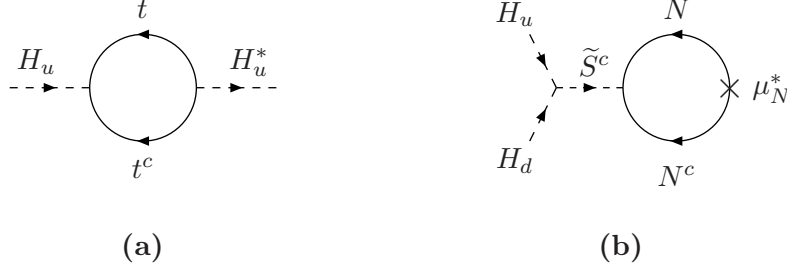


FIG. 2: **(a)** A contribution to the radiatively induced effective Higgs potential by the top quarks in the MSSM. **(b)** A contribution to the radiatively induced effective Higgs potential by the singlets. It is compared with the diagram (a). The basic structures of the loops in (a) and (b) are the same. The trilinear scalar coupling in (b) comes from the cross term of $|\partial W/\partial S|^2$. Radiatively generated mass in the $\{N, N^c\}$ sector is transmitted to the Higgs through the mediation by \tilde{S}^c .

thus, plays the role of a mass generation sector of the Higgs. As seen in FIG. 2-(b), the non-zero mass effect is transmitted to the Higgs through the messenger \tilde{S}^c , which is actually a mediator of the Higgs mass effect. The mass term generated in this way can be meaningful only below the mass scale of \tilde{S}^c ($\approx \mu_S$), because it can be regarded as a local operator below the scale of μ_S .

FIG. 2-(a) and (b) show the typical diagrams for the radiatively generated Higgs potentials by the top quarks in the MSSM and the singlets in our case, respectively. They are compared to each other. Actually, FIG. 2-(b) contributes to the renormalization of $B\mu$ term, while FIG. 2-(a) to the renormalization of the m_2^2 . The basic structures of the loops in the two diagrams are the same. Roughly, the diagram of FIG. 2-(b) is estimated as $(H_u H_d)(y_H \mu_S^*) \frac{1}{\mu_S^2} [n \times \text{Loop}] \mu_N^*$, while FIG. 2-(a) as $H_u y_t [3 \times \text{Loop}] y_t^* H_u^*$, where the “Loop” means the common calculation of the loops in the diagrams.

The radiative correction $\Delta V(H)$ given with Eqs. (12) and (19) expands with the powers of h_u and h_d as follows:

$$\Delta V(H) = \Delta V_{(0,0)} + \Delta V_{(1,1)} h_u h_d + \frac{1}{2!2!} \Delta V_{(2,2)} (h_u h_d)^2 + \cdots, \quad (20)$$

where the coefficients such as $\Delta V_{(0,0)} [\equiv \Delta V(H)|_{h_u=h_d=0}]$, $\Delta V_{(1,1)} [\equiv \partial_{h_u} \partial_{h_d} \Delta V(H)|_{h_u=h_d=0}]$ and $\Delta V_{(2,2)} [\equiv \partial_{h_u}^2 \partial_{h_d}^2 \Delta V(H)|_{h_u=h_d=0}]$ are estimated as

$$\begin{aligned} \Delta V_{(0,0)} &= \frac{n}{16\pi^2} \left[(\mu_N^2 + \tilde{m}_N^2)^2 \left\{ \log \left(\frac{\mu_N^2 + \tilde{m}_N^2}{\Lambda^2} \right) - \frac{3}{2} \right\} - \mu_N^4 \left\{ \log \left(\frac{\mu_N^2}{\Lambda^2} \right) - \frac{3}{2} \right\} \right], \\ \Delta V_{(1,1)} &= \frac{n}{8\pi^2} \left(y_H y_N \frac{\mu_N}{\mu_S} \right) \left[(\mu_N^2 + \tilde{m}_N^2) \left\{ \log \left(\frac{\mu_N^2 + \tilde{m}_N^2}{\Lambda^2} \right) - 1 \right\} - \mu_N^2 \left\{ \log \left(\frac{\mu_N^2}{\Lambda^2} \right) - 1 \right\} \right], \\ \Delta V_{(2,2)} &= \frac{n}{4\pi^2} \left(y_H y_N \frac{\mu_N}{\mu_S} \right)^2 \log \left(\frac{\mu_N^2 + \tilde{m}_N^2}{\mu_N^2} \right). \end{aligned} \quad (21)$$

Note that the coefficients of h_u , h_d , h_u^2 , h_d^2 , h_u^3 , h_d^3 , $h_u^2 h_d$, and $h_u h_d^2$ are all the zeros, and the part of “...” are much suppressed with the higher powers of $(y_H y_N h_{u,d}^2 / \mu_S \mu_N)$. $\Delta V_{(0,0)}$ in Eq. (20) or (21) just adds positive vacuum energy as seen from the first diagram of FIG. 1, which is a result of SUSY breaking.

As the (s-) top quark loops renormalize the soft mass squared of the Higgs, m_2^2 in the MSSM, the diagram of FIG. 2-(b) or $\Delta V_{(1,1)}$ term in Eq. (20) renormalizes the B_μ term in Eq. (18) ($B_\mu \mu \equiv m_3^2$), $m_3^2(\mu) = m_3^2 + \Delta m_3^2$, where

$$\Delta m_3^2 \approx \frac{n}{8\pi^2} \left(y_H y_N \frac{\mu_N}{\mu_S} \right) \left[(\mu_N^2 + \tilde{m}_N^2) \left\{ \log \left(\frac{\mu_N^2 + \tilde{m}_N^2}{\Lambda^2} \right) - 1 \right\} - \mu_N^2 \left\{ \log \left(\frac{\mu_N^2}{\Lambda^2} \right) - 1 \right\} \right]. \quad (22)$$

Below the mass scale of \tilde{S}^c , \tilde{S}^c as well as $\{N, N^c\}$ is decoupled, and the RG running of m_3^2 caused by the interactions between the Higgs and the singlets would become frozen. Thus, we set $\Lambda = \mu_S$ at lower energies.

With the correction Eq. (22), one of the tree-level extremum conditions in the Higgs potential is modified as²

$$-2m_3^2 = (m_1^2 - m_2^2) \tan 2\beta + M_Z^2 \sin 2\beta + 2\Delta m_3^2. \quad (23)$$

In order to avoid a fine-tuning among the parameters, $2\Delta m_3^2$ needs to be comparable with other terms in Eq. (23), when the parameters are chosen for explaining the Higgs mass of 125 GeV. If Δm_3^2 is too large, it should be properly cancelled with other terms, being equated with $M_Z^2 \sin 2\beta$ in Eq. (23). Then, the tuning is roughly estimated with the hierarchy, $M_Z^2 \sin 2\beta / (2\Delta m_3^2)$.

By comparing the quartic term in Eq. (20) with the scalar potential in the NMSSM, $V \supset \lambda^2 |H_u H_d|^2 = \frac{\lambda^2}{4} (h_u h_d)^2$, one can see that $\Delta V_{(2,2)}$ in Eq. (21) plays the role of λ^2 of the NMSSM. Since we saw that the Higgs mass correction to the lightest Higgs mass in the NMSSM is given by $\lambda^2 \times (v_H \sin 2\beta)^2$ in Eq. (7), we can readily get the radiative correction Δm_h^2 in our case:

$$\Delta m_h^2 \approx \frac{n}{4\pi^2} \left(y_H y_N \frac{\mu_N}{\mu_S} \right)^2 (v_H^2 \sin^2 2\beta) \log \left(\frac{\mu_N^2 + \tilde{m}_N^2}{\mu_N^2} \right). \quad (24)$$

Note that μ_S is originated from the propagator of \tilde{S}^c in the diagram, while μ_N from the mass insertion. Thus, the mass term correction by Δm_h^2 can be a *local operator* below the messenger scale μ_S . Since the mass squared of \tilde{S}^c [$> (300 \text{ GeV})^2$] is much heavier than the Higgs mass squared, Δm_h^2 in Eq. (24) indeed can play the role of the Higgs mass correction

² The extremum conditions in the MSSM are $m_1^2 + |\mu|^2 = m_3^2 \tan \beta - \frac{M_Z^2}{2} \cos 2\beta$, and $m_2^2 + |\mu|^2 = m_3^2 \cot \beta + \frac{M_Z^2}{2} \cos 2\beta$ at tree level, which can be recast into $-2m_3^2 = (m_1^2 - m_2^2) \tan 2\beta + M_Z^2 \sin 2\beta$, and $|\mu|^2 = (m_3^2 \sin^2 \beta - m_1^2 \cos^2 \beta) / (\cos 2\beta) - \frac{1}{2} M_Z^2$ [3].

at low energies. By including Eq. (24), the CP even lightest Higgs mass squared is modified as

$$m_h^2 \approx M_Z^2 \cos^2 2\beta - \frac{|A_S - B_S|^2}{\tilde{m}_S^2 + \mu_S^2} (y_H^2 v_H^2 \sin^2 2\beta) + \Delta m_h^2|_{\text{MSSM}} + \Delta m_h^2. \quad (25)$$

With the hierarchy Eq. (11), the classical correction (the second term) is suppressed. Since it turns out to be much smaller than the radiative correction Δm_h^2 , we will neglect it.

As shown in Ref. [12], the Higgs mass of 125 GeV can be explained with Eqs. (25) or (24) in the parameter space,

$$0.2 \lesssim y_H \lesssim 0.7 \quad \text{or} \quad 3 \lesssim \tan\beta \lesssim 10, \quad (26)$$

without considering the mixing effect, if the soft mass of the s-top is around 500 GeV [or $\Delta m_h^2|_{\text{MSSM}} \approx (66 \text{ GeV})^2$]. Thus, even $0.2 \lesssim y_H \lesssim 0.5$ or $3 \lesssim \tan\beta \lesssim 10$, which is the excluded region in the NMSSM, can be still consistent with the 125 GeV Higgs mass, when the radiative correction of the Higgs mass is supported by the MSSM singlet fields.

For the typical three classes, $\mu_S \lesssim \tilde{m}_N \lesssim \mu_N$ (Case A), $\mu_S \lesssim \mu_N \lesssim \tilde{m}_N$ (Case B), and $\mu_S \lesssim \tilde{m}_N \approx \mu_N$ (Case C), the radiative corrections in Eqs. (24) and (22) are approximated as follows:

$$\begin{cases} \Delta m_h^2 \approx \frac{n}{4\pi^2} (v_H^2 \sin^2 2\beta) \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right)^2 \frac{\tilde{m}_N^2}{\mu_N^2} \right] \\ \Delta m_3^2 \approx \frac{n}{4\pi^2} \tilde{m}_N^2 \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right) \log \left(\frac{\mu_N}{\mu_S} \right) \right] \end{cases} \quad \text{for } \tilde{m}_N \lesssim \mu_N \text{ (Case A),} \quad (27)$$

$$\begin{cases} \Delta m_h^2 \approx \frac{n}{4\pi^2} (v_H^2 \sin^2 2\beta) \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right)^2 \log \left(\frac{\tilde{m}_N^2}{\mu_N^2} \right) \right] \\ \Delta m_3^2 \approx \frac{n}{4\pi^2} \tilde{m}_N^2 \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right) \log \left(\frac{\tilde{m}_N}{\mu_S} \right) \right] \end{cases} \quad \text{for } \mu_N \lesssim \tilde{m}_N \text{ (Case B),} \quad (28)$$

$$\begin{cases} \Delta m_h^2 \approx \frac{n}{4\pi^2} (v_H^2 \sin^2 2\beta) \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right)^2 \log 2 \right] \\ \Delta m_3^2 \approx \frac{n}{4\pi^2} \tilde{m}_N^2 \left[\left(y_H y_N \frac{\mu_N}{\mu_S} \right) \left\{ \log \left(2 \frac{\mu_N}{\mu_S} \right) - \frac{1}{2} \right\} \right] \end{cases} \quad \text{for } \mu_N \approx \tilde{m}_N \text{ (Case C).} \quad (29)$$

In order to avoid a serious fine-tuning among soft parameters in Eq. (23), $\Delta m_3^2/v_H^2$ should not be too much larger than the unity. From the above equations, roughly it means $\tilde{m}_N \lesssim 2\pi v_H \approx 1 \text{ TeV}$. Hence, \tilde{m}_N should be quite smaller than 1 TeV. In the next section, we will discuss this issue in more detail.

IV. 125 GEV HIGGS MASS WITH THE LEAST TUNING

For simple presentations, we parameterize the radiative corrections in Eqs. (24) and (22) as follows:

$$F^2 \equiv \frac{\Delta m_h^2}{f^2 v_H^2} = R^2 \log(1 + r^2),$$

$$G \equiv \frac{2\Delta m_3^2}{g\mu_S^2} = R^3 \left[(1 + r^2) \{ \log(1 + r^2) + \log R^2 - 1 \} - \{ \log R^2 - 1 \} \right], \quad (30)$$

where R , r , f^2 and g are defined as

$$R \equiv \frac{\mu_N}{\mu_S}, \quad r \equiv \frac{\tilde{m}_N}{\mu_N}, \quad \text{and}$$

$$f^2 \equiv \frac{n}{4\pi^2} y_H^2 y_N^2 \sin^2 2\beta, \quad g \equiv \frac{n}{4\pi^2} y_H y_N. \quad (31)$$

For the parameters chosen for the explanation of the Higgs mass around 125 GeV, as mentioned above, a smaller Δm_3^2 is more desirable to avoid a fine-tuning among the parameters in Eq. (22). From now on we will explore the conditions under which Δm_3^2 can be minimized for a given Δm_h^2 and other parameters in the model. As seen from Eq. (30), R and r are related each other for a given F . Accordingly, G depends only on r or R for a fixed F . Let us insert F into G , replacing r by R and F . For a given set of $\{\Delta m_h^2, \mu_S^2, f^2, g\}$, thus, G is recast to

$$G = R^3 \left[e^{\frac{F^2}{R^2}} \left(\frac{F^2}{R^2} + \log R^2 - 1 \right) - (\log R^2 - 1) \right]. \quad (32)$$

Provided that F is fixed, one can show that G is minimized at

$$R = \frac{F}{1 + \epsilon_F} \quad (33)$$

where the small parameter ϵ_F is estimated as

$$\epsilon_F \approx \frac{1 - 0.28 \log F^2}{8.87 + 4.31 \log F^2}. \quad (34)$$

$|\epsilon_F|$ is much smaller than the unity in the most parameter range of F : $|\epsilon_F|$ is smaller than 0.3 (0.1) for $0 < |F| < 0.16$ or $0.59 < |F|$ ($0 < |F| < 1.9 \times 10^{-3}$ or $1.08 < |F|$). From Eq. (30), thus, r and G are determined when G is minimized:

$$r^2 \approx 1.72 + 5.44\epsilon_F,$$

$$G \approx F^3 \left[(1.72 + 0.28\epsilon_F) \log F^2 + (1 - \epsilon_F) \right]. \quad (35)$$

For instance, $\epsilon_F \approx 0.05$, $R \approx 1.75$, $r \approx 1.35$, and $G \approx 19.09$ for $F = 1.84$. From Eq. (31), it implies that $\frac{\mu_N}{\mu_S} \approx 1.75$, $\frac{\tilde{m}_N}{\mu_N} \approx 1.35$, and $\Delta m_3^2 \approx (330 \text{ GeV})^2$ e.g. for $|\Delta m_h| = 91 \text{ GeV}$, $\mu_S = 300 \text{ GeV}$, $n = 5$, $y_H y_N = 1$ and $\sin 2\beta = 0.8$.

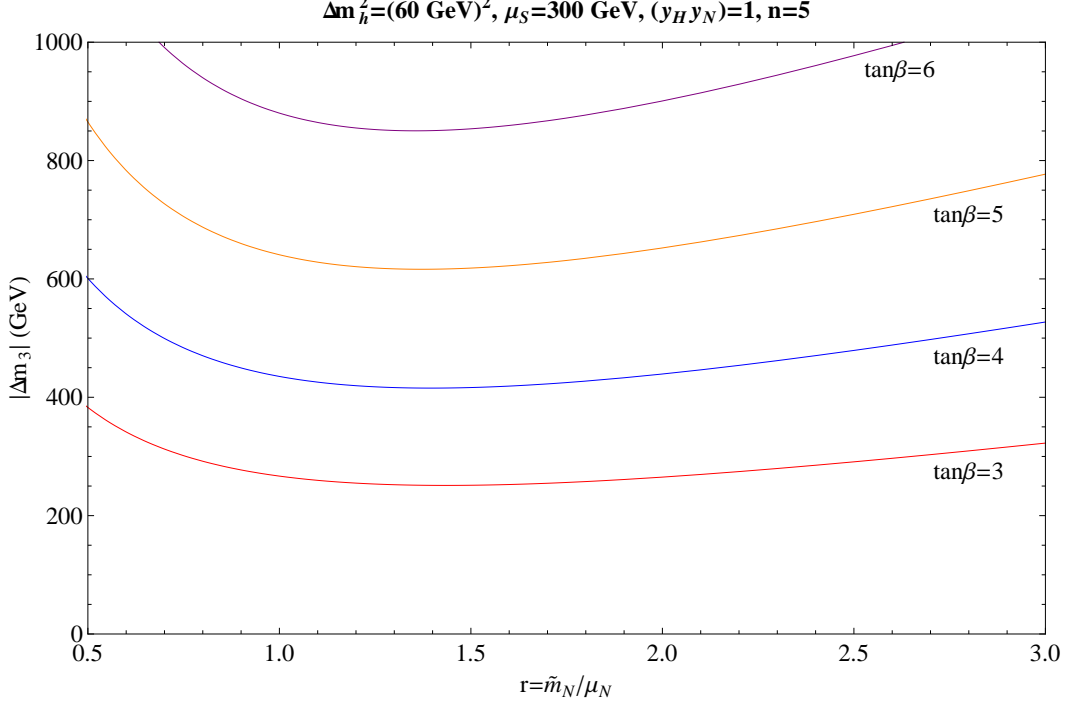


FIG. 3: Radiative correction $|\Delta m_3|$ ($\equiv \sqrt{B_\mu \mu}$) vs. \tilde{m}_N/μ_N for various values of $\tan\beta$. The radiative correction to the Higgs mass Δm_h^2 is set to $(60 \text{ GeV})^2$. Thus, $|\Delta m_h|_{\text{MSSM}} \approx (68, 70, 75, 82)$ GeV for $\tan\beta = (6, 5, 4, 3)$ are assumed to be supplemented from the (s-) top's contributions for the 125 GeV Higgs mass. They correspond to $\tilde{m}_t \approx (530, 590, 780, 1300)$ GeV at two-loop level, when turning off the mixing effect of $(\tilde{t}_L, \tilde{t}_R)$. We fix the other parameters as shown in the figure.

Note that $0.3 < r < 1.8$ for $-0.3 < \epsilon_F < 0.3$ in Eq. (35). We can see that μ_N and \tilde{m}_N need to be comparable to each other in order to minimize Δm_3^2 . However, Δm_3^2 is not much sensitive to \tilde{m}_N/μ_N ($= r$), only if \tilde{m}_N/μ_N is larger than the unity, because r logarithmically depends on the constraint relation associated with F in Eq. (30).

In Eq. (35), G could be further minimized with a small F . Since $\Delta m_h^2 \approx m_h^2 - M_Z^2 \cos^2 2\beta - \Delta m_h^2|_{\text{MSSM}}$, F^2 in Eq. (30) is minimized when $\sin^2 2\beta = 1$ (or $\tan\beta = 1$):

$$F^2 \approx \frac{m_h^2 - M_Z^2 - \Delta m_h^2|_{\text{MSSM}} + M_Z^2 \sin^2 2\beta}{\frac{n}{4\pi^2} (y_H y_N)^2 v_H^2 \sin^2 2\beta} \geq \frac{m_h^2 - \Delta m_h^2|_{\text{MSSM}}}{\frac{n}{4\pi^2} (y_H y_N)^2 v_H^2}, \quad (36)$$

For $n = 5$, $(y_H y_N) = 1$, and $\Delta m_h^2|_{\text{MSSM}} = (66 \text{ GeV})^2$ [which corresponds to $\tilde{m}_t \approx 500$ GeV at two-loop level], thus, the minimum of F^2 is $(1.71)^2$, which gives $G \approx 14.07$ or $\Delta m_3^2 \approx (284 \text{ GeV})^2$. For Δm_3^2 with other parameters, see FIG. 3 and 4.

Let us present the estimations on typical values of Δm_3^2 for the three classes defined in section III, when Δm_h^2 and other parameters given. In Case A, namely, for $\tilde{m}_N \lesssim \mu_N$, we

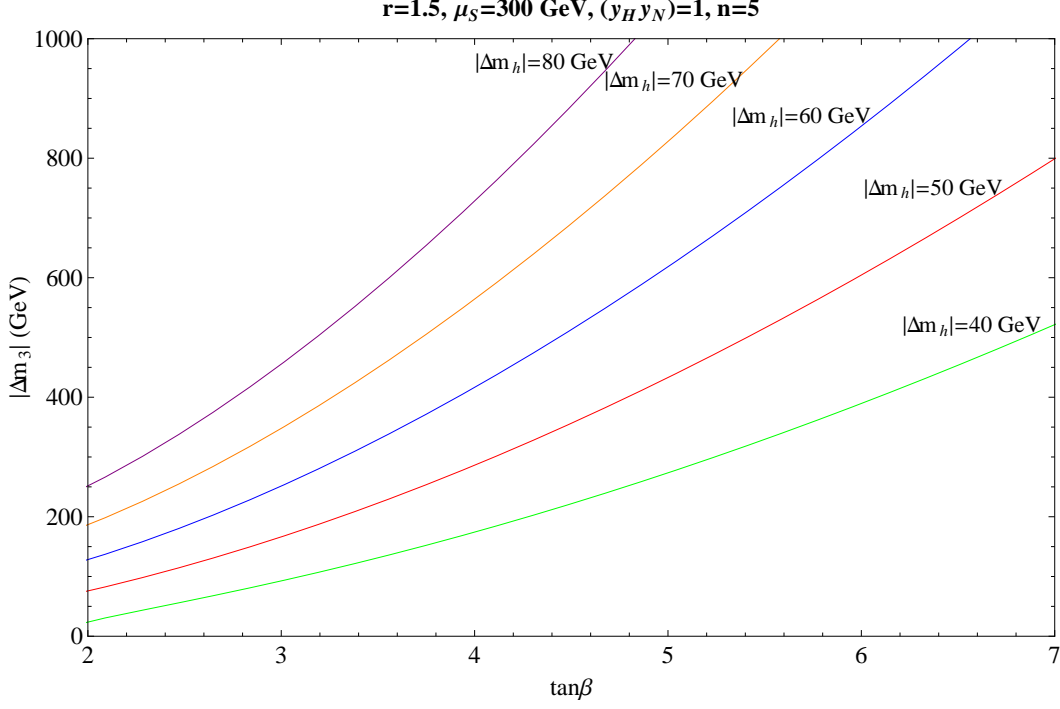


FIG. 4: Radiative correction $|\Delta m_3|$ ($\equiv \sqrt{B_{\mu\mu}}$) vs. $\tan\beta$ for various values of Δm_h^2 around the least tuning points ($\tilde{m}_N/\mu_N \approx 1.5$). $|\Delta m_h| = (80, 70, 60, 50, 40)$ GeV for $\tan\beta = 5$ require the supplements of $|\Delta m_h|_{\text{MSSM}} \approx (47, 61, 70, 78, 83)$ GeV, respectively, by the (s-) top's contributions. They correspond to $\tilde{m}_t \approx (230, 390, 590, 940, 1400)$ GeV at two-loop level, when turning off the mixing effect of $(\tilde{t}_L, \tilde{t}_R)$. The other parameters are fixed as shown in the figure.

have

$$\begin{aligned} \Delta m_3^2 &\approx \mu_S^2 \left[\frac{|\Delta m_h|}{v_H} \right]^3 \frac{g}{f^3 r} \log \left(\frac{|\Delta m_h|}{f v_H r} \right) \\ &\approx (592 \text{ GeV})^2 \left[\frac{\mu_S}{300 \text{ GeV}} \right]^2 \left[\frac{|\Delta m_h|}{90 \text{ GeV}} \right]^3 \left[\frac{1.14}{\sqrt{n}(y_H y_N)^2 \sin^3 2\beta} \right] \left[\frac{\frac{1}{r} \log \left(\frac{|\Delta m_h|}{f v_H r} \right)}{3 \log \left(\frac{3.90}{0.28 \cdot 174} \right)} \right]. \end{aligned} \quad (37)$$

In Case B, i.e. for $\mu_N \lesssim \tilde{m}_N$,

$$\begin{aligned} \Delta m_3^2 &\approx \mu_S^2 \left[\frac{|\Delta m_h|}{v_H} \right]^3 \frac{g r^2}{f^3 (\log r^2)^{3/2}} \log \left(\frac{r |\Delta m_h|}{\sqrt{\log r^2} f v_H} \right) \\ &\approx (499 \text{ GeV})^2 \left[\frac{\mu_S}{300 \text{ GeV}} \right]^2 \left[\frac{|\Delta m_h|}{90 \text{ GeV}} \right]^3 \left[\frac{1.14}{\sqrt{n}(y_H y_N) \sin^3 2\beta} \right] \left[\frac{\frac{r^2}{(\log r^2)^{3/2}} \log \left(\frac{r |\Delta m_h|}{\sqrt{\log r^2} f v_H} \right)}{2.76 \log \left(\frac{3.90}{1.48 \cdot 0.28 \cdot 174} \right)} \right]. \end{aligned} \quad (38)$$

In Case C, i.e. for $\mu_N \approx \tilde{m}_N$,

$$\begin{aligned} \Delta m_3^2 &\approx \mu_S^2 \left[\frac{|\Delta m_h|}{v_H} \right]^3 \frac{g}{f^3 (\log 2)^{3/2}} \left[\log \left(\frac{2|\Delta m_h|}{f v_H \sqrt{\log 2}} \right) - \frac{1}{2} \right] \\ &\approx (342 \text{ GeV})^2 \left[\frac{\mu_S}{300 \text{ GeV}} \right]^2 \left[\frac{|\Delta m_h|}{90 \text{ GeV}} \right]^3 \left[\frac{1.14}{\sqrt{n} (y_H y_N)^2 \sin^3 2\beta} \right] \left[\frac{\log \left(\frac{2|\Delta m_h|}{f v_H \sqrt{\log 2}} \right) - \frac{1}{2}}{0.99} \right]. \end{aligned} \quad (39)$$

V. THE MODEL

The effective superpotential Eq. (10) can be reduced e.g. from the following UV superpotential:

$$\begin{aligned} W_{\text{UV}} &= y_H S H_u H_d + y_N S^c N N^c \\ &+ \frac{f_1}{M_P} \Sigma_1^2 H_u H_d + \frac{f_2}{M_P} \Sigma_2^2 N N^c + \frac{f_3}{M_P} \Sigma_3^2 S S^c \\ &+ \frac{g_1}{M_P} \Sigma_3 \Sigma_1 \bar{\Sigma}_1^2 + \frac{g_2}{M_P} \Sigma_3 \Sigma_2 \bar{\Sigma}_2^2 + \frac{g_3}{M_P} \Sigma_3^2 \bar{\Sigma}_3^2 \end{aligned} \quad (40)$$

where y_H , y_N , f_i , and g_i ($i = 1, 2, 3$) are dimensionless couplings, and M_P denotes the reduced Planck mass ($= 2.4 \times 10^{18} \text{ GeV}$). The superpotential Eq. (40) respects the two continuous global symmetries, $U(1)_R$ and $U(1)_{\text{PQ}}$. The global charges for the superfields in the superpotential Eq. (40) are displayed in TABLE I.

Superfields	H_u	H_d	N	N^c	S	S^c	Σ_1	Σ_2	Σ_3	$\bar{\Sigma}_1$	$\bar{\Sigma}_2$	$\bar{\Sigma}_3$
$U(1)_R$	0	0	0	0	2	2	1	1	-1	1	1	2
$U(1)_{\text{PQ}}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	-1	$\frac{1}{4}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{8}$	$-\frac{1}{4}$

TABLE I: R and Pecci-Quinn charges of the superfields. The MSSM *matter* superfields carry the unit R charges, and also the PQ charges of 1/8. N and N^c are assumed to be proper n -dimensional vector-like representations of a hidden gauge group, under which all the MSSM fields are neutral. Σ s and $\bar{\Sigma}$ s carry some Z_2 charges.

The “A-terms” corresponding to the $g_{1,2,3}$ terms in Eq. (40), the soft mass terms, etc. in the scalar potential admit the VEVs of $\Sigma_{1,2,3}$ and $\bar{\Sigma}_{1,2,3}$ of order $\sqrt{m_{3/2} M_P}$ ($\sim 10^{10} \text{ GeV}$) at the minimum [15]. From the $f_{1,2,3}$ terms in Eq. (40), thus, “ μ ” in the MSSM, and also μ_N and μ_S in Eq. (10) are generated, which are of order $m_{3/2}$ [16]. Due to the VEVs of Σ s and $\bar{\Sigma}$ s, the $U(1)_R$ and $U(1)_{\text{PQ}}$ are completely broken at the intermediate scale.

Σ s and $\bar{\Sigma}$ s in Eq. (40) carry some accidental Z_2 charges. As a result, the domain wall problem would potentially arise. Hence, we assume that such discrete symmetries were already broken before or during inflation such that domain walls were diluted away. If the reheating temperature is lower than 10^9 GeV , the Z_2 breaking vacuum still remains as the minimum of the potential also after inflation [15].

VI. CONCLUSION

In this paper, we proposed a new type of the singlet extension of the MSSM in order to raise the Higgs mass up to 125 GeV with the tuning associated with the light Higgs mass relieved. Apart from the (s-) top quark's contribution, the Higgs mass is radiatively generated in a hidden sector because of the mass splitting of hidden sector fields, and such a effect is transmitted to the Higgs sector through the mediation by the messenger field \tilde{S}^c . Since the Higgs mass is raised by the superpotential parameters, lifting the Higgs mass is quite efficient as in the extra matter scenario. Unlike the extra matter scenario, however, our model is free from the constraint on extra colored particles with order one Yukawa couplings with the Higgs, which is associated with the production and decay rates of the Higgs at the LHC [8].

As shown in our previous paper [12], the parameter space for 125 GeV Higgs mass can be enlarged compared to the original form of the NMSSM, and so even $0.2 \lesssim y_H \lesssim 0.5$ or $3 \lesssim \tan\beta \lesssim 10$, which is excluded region in the NMSSM, can explain the 125 GeV Higgs mass with a relatively light s-top (~ 500 GeV) but without considering the mixing effect. In this paper, we also particularly emphasized that the fine-tuning problem associated with the light Higgs mass can be remarkably mitigated by taking low enough messenger scale (≈ 300 GeV) and light enough mass parameters ($\ll 1$ TeV). We have explored the least tuning condition ($\mu_N \lesssim \tilde{m}_N$), under which even the soft parameters lighter than 500 GeV can explain the Higgs mass of 125 GeV.

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